

**From Boson Condensation to Quark
Deconfinement:
The Many Faces of Neutron Star Interiors**

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February 1, 2008

Presented at the
XXIII School of Theoretical Physics
Ustron, Poland
15–22 September 1999
To be published in the Acta Physica Polonica

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From Boson Condensation to Quark Deconfinement: The Many Faces of Neutron Star Interiors

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Abstract

Gravity compresses the matter in the cores of neutron stars to densities which are significantly higher than the density of ordinary atomic nuclei, thus providing a high-pressure environment in which numerous particle processes – from the generation of new baryonic particles to quark deconfinement to the formation of Boson condensates and H-matter – may compete with each other. There are theoretical suggestions of even more ‘exotic’ processes inside pulsars, such as the formation of absolutely stable strange quark matter, a configuration of matter even more stable than the most stable atomic nucleus, iron. In the latter event, neutron stars would be largely composed of pure quark matter, eventually enveloped in nuclear crust matter. No matter which physical processes are actually realized inside neutron stars, each one leads to fingerprints, some more pronounced than others though, in the observable stellar quantities. This feature combined with the tremendous recent progress in observational radio and X-ray astronomy, renders neutron stars to nearly ideal probes for a wide range of dense matter studies, complementing the quest of the behavior of superdense matter in terrestrial collider experiments.

1 Introduction

Neutron stars are spotted as pulsars by radio telescopes and X-ray satellites. They are more massive (i.e. $\sim 1.5 M_{\odot}$) than our sun but are typically only about ~ 10 kilometers across so that the matter in their centers is compressed to densities that are up to an order of magnitude higher than the density of atomic nuclei. A neutron star, therefore, provides a high-pressure environment in which numerous subatomic particle processes are expected to compete with each other and novel phases of matter – like the quark-gluon plasma being sought at the most powerful terrestrial particle colliders – could exist. In my lecture, I will give an overview of the present status

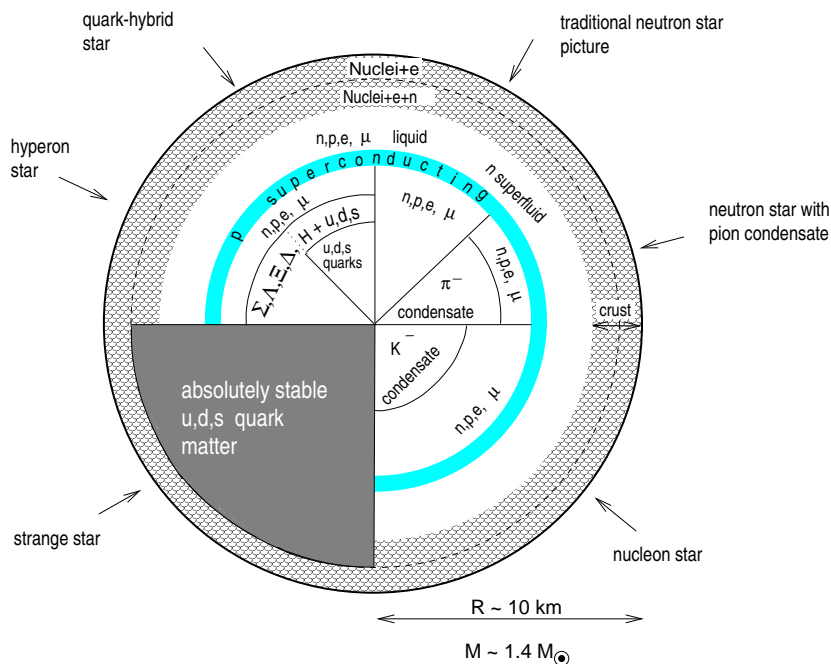


Figure 1: Competing structures and novel phases of subatomic matter predicted by theory to make their appearance in the cores ($R \lesssim 8$ km) of neutron stars [1].

of the research on the many phases of superdense matter in neutron stars, which naturally is to be performed at the interface between nuclear physics, particle physics and Einstein's theory of relativity. Of particular interest will be the existence of quark matter inside neutron stars and the fingerprints by means of which this novel phase of matter could register itself in the observed neutron star data. The detection of such matter in neutron stars would help to clarify how quark matter behaves, and give a boost to theories about the early Universe as well as laboratory searches for the production of quark matter in heavy-ion colliders. Complementary talks on the physics of neutron stars will be given by J. Lattimer and M. Prakash who will be discussing the structure and evolution of neutron stars, and the neutrino interactions in dense matter, respectively.

2 The many faces of neutron stars

From model calculations, it is known that neutron stars are far from being composed of only neutrons but instead may possess rather complex interior structures, as established in model calculations performed over the years. Figure 1 gives an overview of those structures that are currently most vividly discussed in the literature (for an overview, see [1]). No matter which physical structures are actually realized inside neutron stars, each one leads to fingerprints, some more pronounced than others though, in the equation of state (i.e. pressure versus density) associated with these

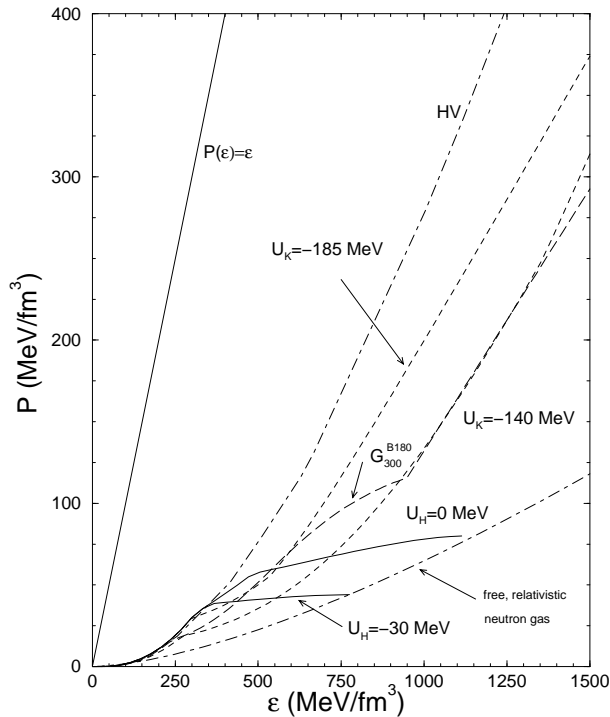


Figure 2: Models for the equation of state (EoS) of ‘neutron’ star matter [1, 2, 3].

phases of matter. This becomes very evident from Fig. 2, which shows a collection of competing models for the equation of state of neutron star matter [1].

2.1 Hyperon star

Only in the most primitive conception, a neutron star is constituted from neutrons. At a more accurate representation, a neutron star will contain neutrons (n) and a small number of protons (p) whose charge is balanced by leptons (e^- and μ^-). The interactions among the nucleons can be treated in the framework of either Schrodinger-based theories [4, 5], the semiclassical Thomas-Fermi method [6], or relativistic nuclear field theories solved at the mean-field level [1, 7] or beyond [1, 8, 9]. The coupling constants of the theory must reproduce the bulk properties of nuclear matter at saturation density, $\rho_0 = 0.16 \text{ fm}^{-3}$ (energy density of $\epsilon_0 = 140 \text{ MeV/fm}^3$). These are the binding energy E/A , effective nucleon mass m_N^*/m_N , incompressibility K , and the symmetry energy a_s whose respective values are

$$E/A = -16.0 \text{ MeV}, \quad m_N^*/m_N = 0.79, \quad K \simeq 265 \text{ MeV}, \quad a_s = 32.5 \text{ MeV}. \quad (1)$$

Of the five, the value for the incompressibility of nuclear matter carries some uncertainty. Its value is currently believed to lie in the range between about 180 and 300 MeV. At the densities in the interior of neutron stars, the neutron chemical potential, μ^n , exceeds the mass (modified by interactions) of various members of the

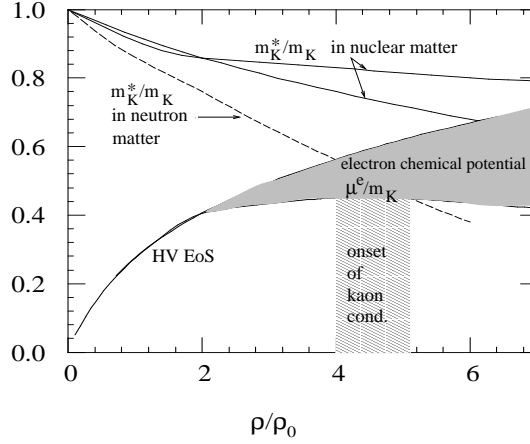


Figure 3: Effective kaon mass in nuclear [12] and neutron star [13] matter.

baryon octet [10]. So in addition to nucleons and electrons, neutron stars may be expected to have populations of hyperons, i.e. Σ , Λ , Ξ and eventually of Δ 's.

2.2 Nucleon stars

Once the reaction

$$e^- \rightarrow K^- + \nu \quad (2)$$

becomes possible in a neutron star, it becomes energetically advantageous for the star to replace the fermionic electrons with the bosonic K^- mesons. Whether or not this actually happens depends on the mass of the K^- in dense matter. A handle on this is provided by the K^- kinetic energy spectra extracted from Ni+Ni collisions at SIS energies, measured by the KaoS collaboration at GSI [11]. An analysis of the KaoS data shows that the attraction from nuclear matter may bring the K^- mass down to $m_{K^-}^* \simeq 200$ MeV at $\rho \sim 3\rho_0$. For neutron-rich matter, the relation

$$m_{K^-}^*(\rho) \simeq m_{K^-} \left(1 - 0.2 \frac{\rho}{\rho_0} \right) \quad (3)$$

was established [14, 15, 16, 17], with $m_K = 495$ MeV the K^- vacuum mass. Values around $m_{K^-}^* \simeq 200$ MeV lie in the vicinity of the electron chemical potential, μ^e , in neutron star matter [1, 10] so that the threshold condition for the onset of K^- condensation, $\mu^e = m_{K^-}^*$, which follows from Eq. (2), could be fulfilled in the cores of neutron stars. The situation is illustrated graphically in Fig. 3. Equation (2) is followed by

$$n + e^- \rightarrow p + K^- + \nu, \quad (4)$$

with the neutrinos leaving the star. By this conversion the nucleons in the cores of newly formed neutron stars can become half neutrons and half protons, which lowers the energy per baryon of the matter [16]. The relatively isospin symmetric

composition achieved in this way resembles the one of atomic nuclei, which are made up of roughly equal numbers of neutrons and protons. Neutron stars are therefore referred to, in this picture, as nucleon stars. The maximal possible mass of this type of star, where Eq. (4) has gone to completion, has been calculated to be between about $1.5 M_{\odot}$ [18] and $1.8 M_{\odot}$ [2]. Based on the former mass value, Brown et al. studied in a recent paper the formation and evolution of black holes in the Galaxy [19].

Meson condensates can soften the equation of state (EoS) considerably. As a consequence, neutron stars with K^- condensates can be rather dense and, therefore, have radii smaller than neutron stars without condensates, i.e. $R \lesssim 10$ km. An interesting candidate of such a small-radius object may be the nearby neutron star RXJ 185 635–3754 [20] whose radius could be as small as ~ 7 km [21].

2.3 H-dibaryons

A novel particle that could make its appearance in the center of a neutron star is the so-called H-dibaryon, a doubly strange six-quark composite with spin and isospin zero, and baryon number two [22]. Since its first prediction in 1977, the H-dibaryon has been the subject of many theoretical and experimental studies as a possible candidate for a strongly bound exotic state. In neutron stars, which may contain a significant fraction of Λ hyperons, the Λ 's could combine to form H-dibaryons, which could give way to the formation of H-matter at densities somewhere between $3\epsilon_0$ [3] and $6\epsilon_0$ [23, 24], depending on the in-medium properties of the H-dibaryon. H-matter could thus exist in the cores of moderately dense neutron stars. In [3] it was pointed out that H-dibaryons with a vacuum mass of about 2.2 GeV and a moderately attractive potential in the medium of about -30 MeV could go into a Bose condensate in the cores of neutron stars if the limiting star mass is about that of the Hulse–Taylor pulsar PSR 1913+16, $M = 1.444 M_{\odot}$. Conversely, if the medium potential were moderately repulsive, around $+30$ MeV, the formation of H-dibaryons may only take place in heavier neutron stars of mass $M \gtrsim 1.6 M_{\odot}$. If formed, however, H-matter may not remain dormant in neutron stars but, because of its instability against compression could trigger the conversion of neutron stars into hypothetical strange stars [24, 25, 26].

2.4 Quark deconfinement in neutron stars

It has been suggested already back in the 1970's by a number of researchers [27, 28, 29, 30, 31, 32] that, because of the extreme densities reached in the cores of neutron stars, neutrons protons plus the heavier constituents ($\Sigma, \Lambda, \Xi, \Delta$) can melt, creating the quark-gluon plasma state being sought at the most powerful terrestrial heavy-ion colliders at CERN [i.e. experiments NA35 NA44, NA45, CERES, and NA50; in a few years at the LHC by the ALICE experiment] and RHIC. At present one does not know from experiment at what density the expected phase transition to quark matter occurs, and one has no conclusive guide yet from lattice QCD simulations. From simple geometrical considerations it follows that nuclei begin to touch each other at

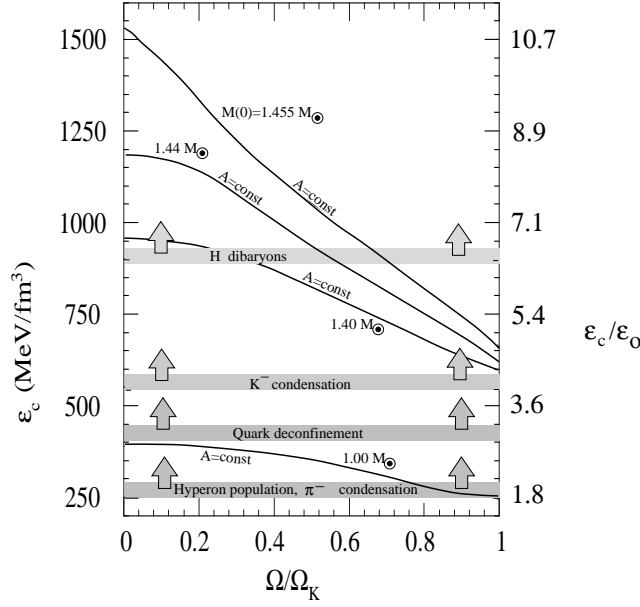


Figure 4: Central density versus rotational frequency of sample neutron stars of constant baryon number, A . Estimated threshold densities of various novel phases of superdense matter are indicated. $M(0)$ is the non-rotating star mass, Ω_K [cf. Eq. (9)] stands for the Kepler frequency which terminates stable rotation [33].

densities of $\sim (4\pi r_N^3/3)^{-1} \simeq 0.24 \text{ fm}^{-3}$, which, for a characteristic nucleon radius of $r_N \sim 1 \text{ fm}$, is less than twice the baryon number density ρ_0 of ordinary nuclear matter [7]. Above this density, therefore, it appears plausible that the nuclear boundaries of hadrons dissolve so that the formerly confined quarks now populate free states outside of the hadrons. Depending on rotational frequency and stellar mass, densities as large as two to three times ρ_0 are easily surpassed in the cores of neutron stars, as can be seen from Fig. 4, so that the neutrons and protons in the centers of neutron stars may have been broken up into their constituent quarks by gravity [33]. More than that, since the mass of the strange quark is only $m_s \sim 150 \text{ MeV}$, high-energetic up and down quarks are expected to readily transform to strange quarks at about the same density at which up and down quark deconfinement sets in [34, 35]. Three flavor quark matter could thus exist as a permanent component of matter in the centers of neutron stars [1, 7, 33, 36]. As we shall see in Section 5, radio astronomers may be able to spot evidence for the existence of this novel phase of matter in the timing structure of pulsar spin-down [33, 36].

2.5 Diquark condensation and color superconductivity

Very recently it was discovered that instantons may cause strong correlations between up and down quarks, which could give way to the existence of colored diquark pairs in superdense matter [37, 38]. These pairs could form a Bose condensate in cold

($T < 50$ MeV) and dense ($\rho > 3\rho_0$) quark matter. Moreover, the condensate ought to exhibit color superconductivity [38, 39]. Both the magnitude of the gap and the critical temperature associated with the color superconductive phase were estimated to be on the order of ~ 100 MeV [38, 39]! The implications of such tremendous gaps for the magnetic fields of pulsars and their thermal evolution were explored in [40, 41] and [42], respectively. Future theoretical studies of QCD at finite baryon number density may reveal to which extent these newly established features of quark matter will have their correspondence in a more complete treatment of QCD at finite baryon number density. They may also shed light on whether or not a diquark condensate could possibly alter the equation of state of neutron star matter sufficiently strongly so that one may expect distinguishing features in the global properties of neutron stars, too.

2.6 Absolutely stable quark matter: the material of strange stars

So far we have assumed that quark matter forms a state of matter higher in energy than atomic nuclei. This most ‘plausible’ assumption, however, may be quite deceiving [43, 44, 45] because for a collection of more than a few hundred u , d , s quarks, the energy per baryon (E/A) of quark matter can be just as well *below* the energy of the most stable atomic nucleus, ^{56}Fe , whose energy per baryon number is $M(^{56}\text{Fe})c^2/56 = 930.4$ MeV, with $M(^{56}\text{Fe})$ the mass of the ^{56}Fe atom. A simple estimate indicates that for strange quark matter $E/A = 4B\pi^2/\mu^3$, so that bag constants of $B = 57.5$ MeV/fm³ (i.e. $B^{1/4} = 145$ MeV) and $B = 85.3$ MeV/fm³ ($B^{1/4} = 160$ MeV) place the energy per baryon of such matter at $E/A = 829$ MeV and 915 MeV, respectively [1, 46, 47, 48]. Obviously, these values correspond to quark matter which is absolutely bound with respect to ^{56}Fe . In this event the ground state of the strong interaction would be strange quark matter (strange matter), made up of u , d , s quarks, instead of nuclear matter. This finding is one of the most startling predictions of modern physics with far-reaching implications for neutron stars, for all hadronic stellar configurations in Fig. 1 would then be only metastable with respect to a stars made up of absolutely stable 3-flavor strange quark matter [1, 44, 49, 50, 51]. If this is indeed the case, and if it is possible for neutron matter to tunnel to quark matter in at least some neutron stars, then it appears likely that all neutron stars would in fact be strange stars [46, 47, 48, 52, 53]. In sharp contrast to the other stars in Fig. 1, which are made up of hadronic matter, possibly in phase equilibrium with quarks, strange stars consist nearly entirely of pure 3-flavor quark matter, eventually enveloped in a thin nuclear crust whose density is less than neutron drip (4×10^{11} g/cm³) [54].

The hypothetical, absolute stability of strange matter gives way to a variety of novel stable strange matter objects which stretch from strangelets at the small baryon number end, $A \sim 10^2$, to strange stars at the high end, $A \sim 10^{57}$, where strange matter becomes unstable against gravitational collapse [35]. The strange counterparts of ordinary atomic nuclei are the strange nuggets, searched for in high-energy collisions

at Brookhaven (e.g. E858, E864, E878, E882-B, E886, E896-A), CERN (Newmass experiment NA52), balloon-borne experiments (CRASH), and terrestrial experiments (e.g. HADRON) [1]. (For a recent review, see [55].) Strange stars should possess properties that may allow one to distinguish them from their ‘conventional’ counterparts [1, 7, 49, 56, 57, 58, 59].

3 Stellar structure equations

Since neutron stars are objects of highly compressed matter, the geometry of space-time is changed considerably from flat space. Neutron star models are thus to be constructed from Einstein’s field equations of general relativity ($\mu, \nu=0,1,2,3$),

$$G^{\mu\nu} \equiv R^{\mu\nu} - \frac{1}{2}g^{\mu\nu}R = 8\pi T^{\mu\nu}(\epsilon, P(\epsilon)), \quad (5)$$

which couples Einstein’s curvature tensor, $G^{\mu\nu}$, to the energy–momentum density tensor, $T^{\mu\nu}$, of the stellar matter. The quantities $g^{\mu\nu}$ and R in (5) denote the metric tensor and the Ricci scalar [1]. Theories of superdense matter enter in Eq. (5) via $T^{\mu\nu}$, which contains the equation of state, $P(\epsilon)$, of the stellar matter. It is derivable from a given stellar-matter Lagrangian $\mathcal{L}(\{\phi\})$ [1]. In general, \mathcal{L} is a complicated function of the numerous hadron and quark fields, collectively written as $\{\phi\}$, that acquire finite amplitudes up to the highest densities reached in the cores of compact stars. According to what has been said in Section 2, plausible candidates for ϕ are the charged states of the SU(3) baryon octet, $p, n, \Sigma, \Lambda, \Xi$ [10], the charged states of the Δ [60, 61], π^- [62] and K^- [2, 14, 15, 16, 17] mesons, as well as the u, d, s quark fields [35]. The conditions of chemical equilibrium and electric charge neutrality of the stellar matter require the presence of leptons too, in which case $\phi = e^-, \mu^-$. Models for the equation of state then follow according to the scheme (for details, see [1])

$$\frac{\partial \mathcal{L}(\{\phi\})}{\partial \phi} - \partial_\mu \frac{\partial \mathcal{L}(\{\phi\})}{\partial (\partial_\mu \phi)} = 0 \quad \Rightarrow \quad P(\epsilon). \quad (6)$$

In general, Eqs. (5) and (6) were to be solved simultaneously since the particles move in curved spacetime whose geometry, determined by Einstein’s field equations, is coupled to the total mass energy ϵ of the matter. In the case of neutron stars, however, the long-range gravitational forces can be cleanly separated from the short-range forces, so that Eqs. (5) and (6) constitute two decouple problems.

3.1 Non-rotating stars

The structure equation of spherical neutron stars has been derived from Einstein’s equation (5) first by Tolman [63], and Oppenheimer and Volkoff [64]. It reads

$$\frac{dP}{dr} = - \frac{\epsilon(r)m(r)}{r^2} \frac{(1 + P(r)/\epsilon(r))(1 + 4\pi r^3 P(r)/m(r))}{1 - 2m(r)/r}, \quad (7)$$

and is known in the literature as the Tolman-Oppenheimer-Volkoff equation, applicable to stellar configurations in hydrostatic equilibrium. We use units for which the gravitational constant and velocity of light are $G = c = 1$ so that the mass of the sun is $M_\odot = 1.47$ km. The mass $m(r)$ contained in a sphere of radius r is given by $m(r) = 4\pi \int_0^r r'^2 \epsilon(r') dr'$. Hence the star's total mass follows as $M \equiv m(R)$.

3.2 Rotating stars

The stellar equations describing rotating compact stars are considerably more complicated than those of non-rotating compact stars [1]. These complications have their cause in the deformation of rotating stars plus the general relativistic effect of the dragging of local inertial frames. This reflects itself in a metric of the form [1, 65]

$$ds^2 = -e^{2\nu} dt^2 + e^{2\psi} (d\phi - \omega dt)^2 + e^{2\mu} d\theta^2 + e^{2\lambda} dr^2, \quad (8)$$

where each metric function, i.e. ν , ψ , μ and λ , depends on the radial coordinate r , polar angle θ , and implicitly on the star's angular velocity Ω . The quantity ω denotes the angular velocity of the local inertial frames, which are dragged along in the direction of the star's rotation. This frequency too depends on r , θ and Ω . Of particular interest is the relative frame dragging frequency $\bar{\omega}$ defined as $\bar{\omega}(r, \theta, \Omega) \equiv \Omega - \omega(r, \theta, \Omega)$, which typically increases from about 15% at the surface to about 60% at the center of a neutron star that rotates at its Kepler frequency [1, 33]. The Kepler frequency, Ω_K , is the maximum frequency a star can have before mass loss (mass shedding) at the equator sets in. It sets an absolute upper limit on stable rapid rotation. In classical mechanics the expression for Ω_K , determined by the equality between centrifuge and gravity, is readily obtained as $\Omega_K = \sqrt{M/R^3}$. Its general relativistic counterpart is given by [1, 65]

$$\Omega_K = \omega + \frac{\omega'}{2\psi'} + e^{\nu-\psi} \sqrt{\frac{\nu'}{\psi'} + \left(\frac{\omega'}{2\psi'} e^{\psi-\nu}\right)^2}, \quad P_K \equiv \frac{2\pi}{\Omega_K}. \quad (9)$$

The primes denote derivatives with respect to the Schwarzschild radial coordinate. Equation (9) is to be evaluated self-consistently together with Einstein's field equations (5) for a given model for the equation of state.

Figure 5 shows Ω_K as a function of rotating star mass. The rectangle indicates both the approximate range of observed neutron star masses as well as the observed rotational periods which, currently, are $P \geq 1.6$ ms. One sees that all pulsars so far observed rotate below the mass shedding frequency and so can be interpreted as rotating neutron stars. Half-millisecond periods or even smaller ones are completely excluded for neutron stars of mass $1.4 M_\odot$ [65, 66, 67] for these equations of state [1]. The situation appears to be very different for neutron stars made up of self-bound strange quark matter, the so-called strange stars introduced in Section 2.6. Such stars appear to withstand stable rotation against mass shedding down to rotational periods in the half-millisecond regime or even less [68]. As a consequence, the possible future discovery of a single sub-millisecond pulsar, say 0.5 ms, could give a strong hint that, firstly, strange stars actually exist and, secondly, the deconfined, self-bound phase of

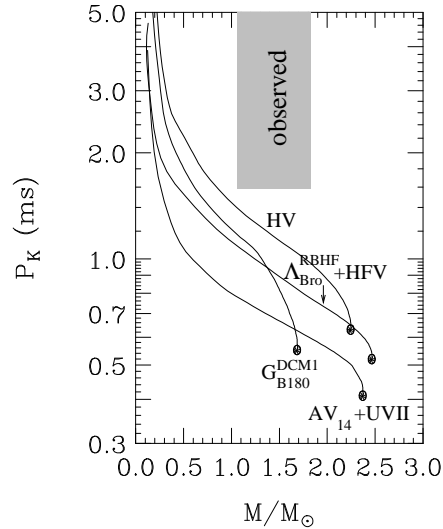


Figure 5: Onset of mass shedding from rapidly spinning neutron stars, computed for a sample of equations of state [1]. The Kepler period is defined in Eq. (9).

3-flavor strange quark matter is the true ground state of the strong interaction rather than nuclear matter. This conclusion is strengthened by the finding of [69] that young strange stars appear not to be subject to the recently discovered r -mode instability, which would slow down hot neutron stars to periods of several milliseconds via the emission of gravitational radiation within a year after birth.

4 Mass constraints from QPOs in LMXBs

Figure 6 exhibits the mass of non-rotating neutron stars as a function of star radius. Each stellar sequence is shown up to the maximum-mass star, indicated by tick marks. Stars beyond the mass peak are unstable against radial oscillations and would collapse to black holes. Evidently, all equations of state can accommodate a neutron star as heavy as the Hulse-Taylor pulsar, whose mass is very precisely known to be $M(\text{PSR } 1913 + 16) = 1.444 M_\odot$. Rather heavy neutron stars, on the other hand, of mass $M \sim 2 M_\odot$ can only be obtained for equations of state that exhibit a rather stiff behavior at supernuclear densities (cf. Fig. 2).

Knowledge of the maximum-mass value is of great importance for two reasons. Firstly, about 20 neutron star masses are presently known [71], and the largest of these imposes a lower bound on the maximum mass of a theoretical model. The current lower bound on the maximum-mass is about $1.56 M_\odot$ (i.e. neutron star 4U 0900–40), which, if firmly established, would indicate that the equation of state of superdense matter will be rather soft at supernuclear densities. This would change if the maximum-mass should be closer to the upper bound of $1.98 M_\odot$. Indications for the possible existence of such heavy neutron stars, with masses around $2 M_\odot$,

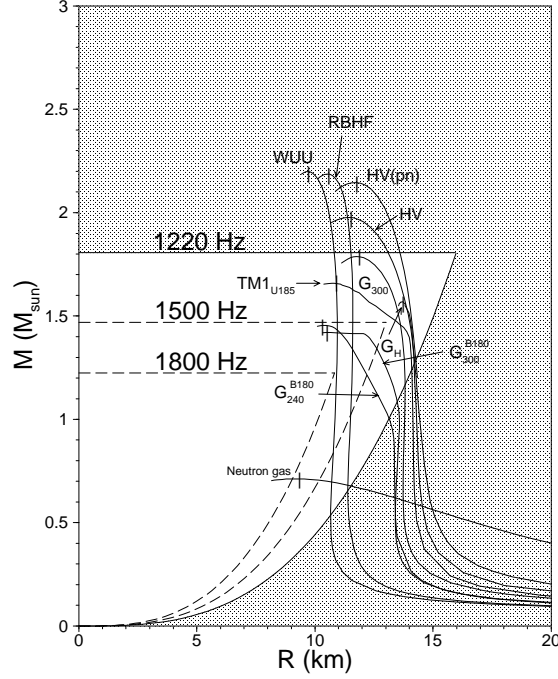


Figure 6: Neutron star mass versus radius for several equations of state shown in Figure 2. Constraints derived from the observation of a QPO frequency of 1220 Hz in neutron star 4U 1636–536 are indicated by the shaded area [70].

may come from the observation of quasi-periodic oscillations (QPOs) in luminosity in low-mass X-ray binaries (LMXBs) [70, 72, 73, 74]. If confirmed, a significant fraction of equations of state presently discussed in the literature [1, 75] could be ruled out. Because of the significant stiffness of the EoS required by heavy neutron stars, all the phase transitions reviewed in Section 2, which imply generally a softening of the EoS rather than a stiffening, appear to be unfavored if not completely ruled out. In this event neutron star matter were likely to be made up of chemically equilibrated nucleons only [75]. Rapid neutron star rotation increases the non-rotating maximum mass value by at most 25% [56, 65, 76, 77, 78, 79], so that even extremely rapidly spinning neutron stars can not have masses significantly above $2.5 M_{\odot}$, as can be seen from Fig. 5. The second reason is that the maximum mass can be useful in identifying black hole candidates [80, 81, 82]. For example, if the mass of a compact companion of an optical star is determined to exceed the maximum mass of a neutron star, it must be a black hole. Because the maximum mass of stable neutron stars in our theory is $2.2 M_{\odot}$, compact companions being more massive than that value are predicted to be black holes. An example of such an object is the non-pulsating X-ray binary Cyg X–1, as its mass lies in the range $9 \lesssim M/M_{\odot} \lesssim 15$.

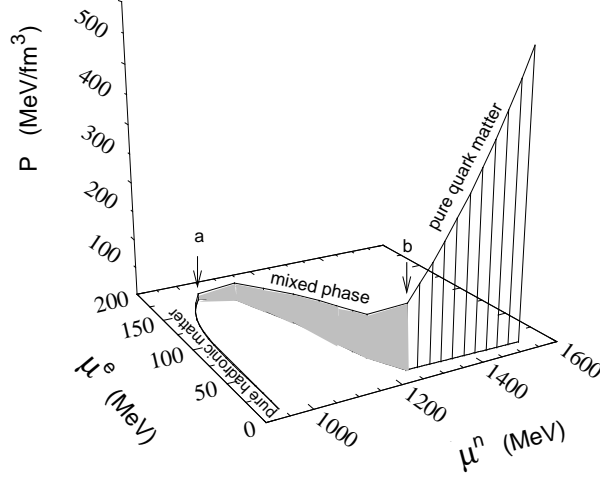


Figure 7: EoS of neutron star matter which accounts for quark deconfinement [1, 33].

5 Evidence of quark matter in neutron stars

As pointed out in Section 2.4, the possibility of quark deconfinement in the cores of neutron stars has already been suggested in the 1970's. However until recently no stringent observational signal has ever been proposed. This is so because whether or not the quark-hadron phase transition occurs in neutron stars makes only little difference to their static properties such as the range of possible masses, radii, or even their limiting rotational periods. This, however, turns out to be strikingly different for the timing structure of rotating neutron stars (i.e. pulsars) that develop quark matter cores in the course of spin-down, as I shall describe in this section.

A model of an EoS which accounts for quark deconfinement in neutron star matter is shown in Fig. 7. One reads off that the transition of confined hadronic matter to quark matter sets in at about twice nuclear matter density (arrow labeled 'a'), which leads to a pronounced softening of the EoS. Pure quark matter exists at densities $\gtrsim 7\epsilon_0$ (arrow labeled 'b').

These structures manifest themselves inside rotating neutron stars as shown in Figs. 8 and 9 (non-rotating star mass is $M = 1.42 M_\odot$). The stars' baryon number is kept constant during spin-down from the Kepler frequency to zero rotation, which describes the temporal evolution of an isolated rotating neutron star. Depending on frequency, the star has an inner sphere of pure quark matter (labeled 'quark' in Fig. 8) surrounded by a few kilometers thick shell of mixed phase of hadronic and quark matter arranged in a Coulomb lattice structure, and this surrounded by a thin shell of hadronic liquid, itself with a thin crust of heavy ions. The lattice structure of varying geometry may have dramatic effects on pulsar observables including transport

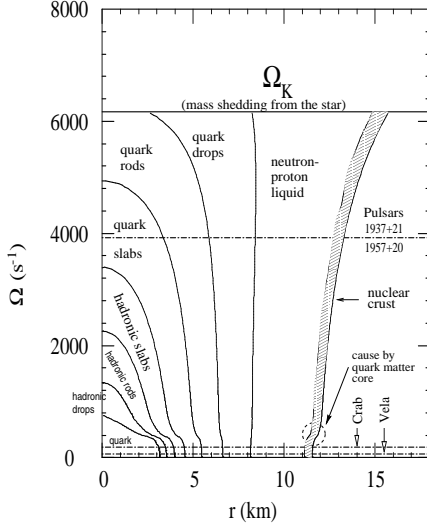


Figure 8: Frequency dependence of quark-hadron structure in equatorial direction of a neutron star [1, 33].

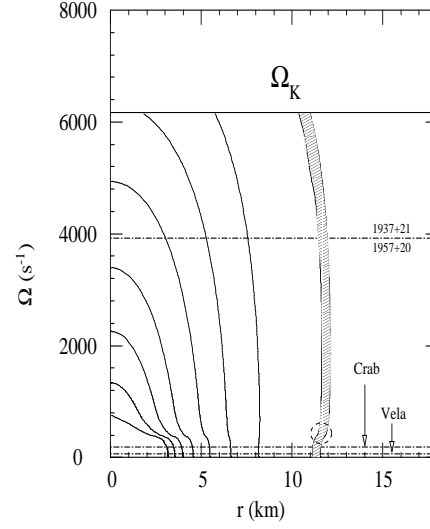


Figure 9: Same as Fig. 8, but in star's polar direction [1, 33].

properties and the theory of glitches [7]. As the star spins up it becomes more and more deformed, and the central density decreases. For some rotating neutron stars the mass and initial rotational frequency may be just such that the central density rises from below to above the critical density for dissolution of baryons into their quark constituents. This is accompanied by a sudden shrinkage of the neutron star, which occurs for the present model star at $\Omega \sim 400 \text{ s}^{-1}$. This effects the star's moment of inertia, I , dramatically. Changes of I , in turn, reflect themselves in the braking index, n , of a pulsar, which is given by ($I' \equiv dI/d\Omega$, $I'' \equiv d^2I/d\Omega^2$) [1, 36],

$$n(\Omega) \equiv \frac{\Omega \ddot{\Omega}}{\dot{\Omega}^2} = 3 - \frac{3 I' \Omega + I'' \Omega^2}{2 I + I' \Omega}. \quad (10)$$

One sees that the braking index depends explicitly and implicitly on Ω . The right side reduces to the canonical constant $n = 3$ only if I is independent of frequency. Figure 10 shows the variation of n with frequency for the star discussed just above (i.e. Figs. 8 and 9) as well as the sample star of [36]. Because of the change in the moment of inertia driven by the transition into quark matter, the braking index deviates dramatically from 3 at the transition frequencies. For the star of Figs. 8 and 9, this occurs at $\Omega \sim 400 \text{ s}^{-1}$ (dashed curve). The solid curve shows the case of a strong deconfinement transition at $\Omega \sim 1370 \text{ s}^{-1}$ discussed in [36]. Continuously connected intermediate cases are possible too and were discussed in [83]. Such dramatic anomalies in $n(\Omega)$ are not known to be exhibited by conventional neutron stars, because their moments of inertia increase smoothly with Ω [1]. The radio astronomical observation of such an anomaly may thus be interpreted as a signal for the development of quark-matter cores in the centers of pulsars!

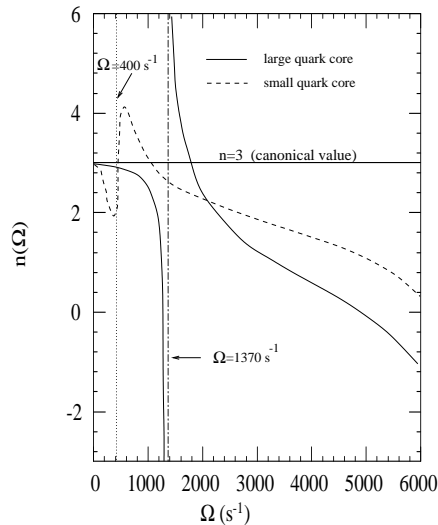


Figure 10: Braking index as a function of rotational frequency [33].

As a very important subject on this issue, we estimate the duration over which the braking index is anomalous. It can be estimated from $\Delta T \simeq -\frac{\Delta\Omega}{\Omega} = \frac{\Delta P}{P}$, where $\Delta\Omega$ is the frequency interval of the anomaly. For a millisecond pulsar whose period derivative is typically $\dot{P} \simeq 10^{-19}$, one finds $\Delta T \simeq 10^8$ years. The dipole age of such pulsars is about 10^9 years. So as a rough estimate we may expect that about 10% of the ~ 25 solitary millisecond pulsars presently known, are in the transition epoch and so could be signaling the ongoing process of quark deconfinement, complementing the searches for the quark-gluon plasma state at the terrestrial heavy-ion colliders.

6 Cooling of neutron stars

The predominant cooling mechanism of hot (temperatures of several $\sim 10^{10}$ K) newly formed neutron stars immediately after formation is neutrino emission, with an initial cooling time scale of seconds. Already a few minutes after birth, the internal neutron star temperature drops to $\sim 10^9$ K [84]. Photon emission overtakes neutrinos only when the internal temperature has fallen to $\sim 10^8$ K, with a corresponding surface temperature roughly two orders of magnitude smaller. Neutrino cooling dominates for at least the first 10^3 years, and typically for much longer in standard cooling (modified Urca) calculations. Being sensitive to the adopted nuclear equation of state, the neutron star mass, the assumed magnetic field strength, the possible existence of superfluidity, meson condensates, quark matter etc., theoretical cooling calculations, as summarized in Fig. 11, too provide most valuable information about the interior matter and neutron star structure. The stellar cooling tracks in Fig. 11 are computed for a broad collection of equations of state [1], including those shown in Fig. 2, which account for the effects mentioned just above. Knowing the thermal evolution of a neutron star also yields information about such temperature-sensitive properties as

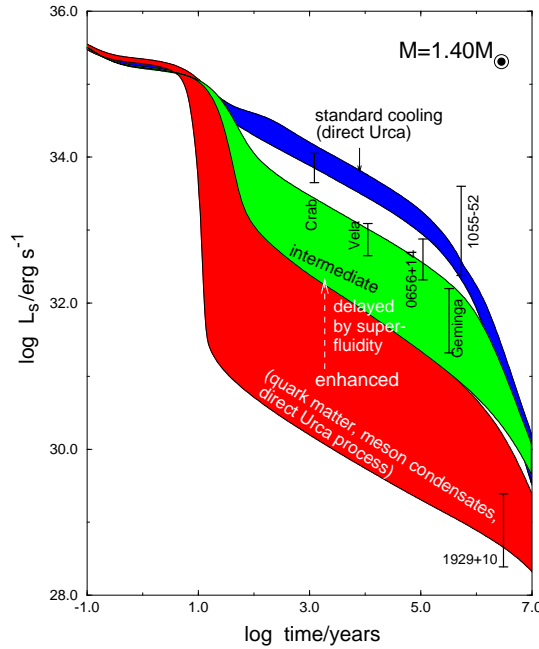


Figure 11: Cooling behavior of a $1.4 M_{\odot}$ neutron star based on competing assumptions about the behavior of superdense matter. Three distinct cooling scenarios, referred to as ‘standard’, ‘intermediate’, and ‘enhanced’ (for details, see [1]), can be distinguished. The band-like structures reflect the uncertainties inherent in the underlying EoS (cf. Fig. 2).

transport coefficients, transition to superfluid states, crust solidification, and internal pulsar heating mechanisms such as frictional dissipation at the crust-superfluid interfaces (for a recent overview, see [85]).

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